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QCD Aspects of Exclusive B Decays*

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Abstract

Exclusive B decays can be factorized as convolutions of hard scattering amplitudes involving the weak interaction with universal hadron distribution amplitudes, thus providing a new QCD-based phenomenology. In addition, semi-leptonic decay amplitudes can be computed exactly in terms of the diagonal and off-diagonal $\Delta n = 2$ overlap of hadronic light-cone wavefunctions. I review these formalisms and the essential QCD ingredients. A canonical form of the light-cone wavefunctions, valid at low values of the transverse momenta, is presented. The existence of intrinsic charm Fock states in the B meson wavefunction can enhance the production of final states of B-decay with three charmed quarks, such as $B \to J/\psi D\pi$, as well as lead to the breakdown of the CKM hierarchy.

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1 Introduction

Remarkable progress has recently been made applying perturbative QCD methods to the exclusive two-body hadronic decays of B mesons. Two groups, Beneke, Buchalla, Neubert, and Sachrajda [1], and Keum, Li, and Sanda [2] have proven factorization theorems which allow the rigorous computation of certain types of exclusive B decay amplitudes in terms of the distribution amplitudes of the initial-state B meson and the final-state hadrons. These new analyses allow one to understand the hadronic physics of heavy hadron decays from a fundamental perspective. There have been many applications of the PQCD formalism, including many new results presented at this conference.

As an example, consider the three representative contributions to the decay of a B meson to meson pairs illustrated in Fig. 1. In Fig. 1(a) the weak interaction effective operator $\mathcal{O}$ produces a $q\overline{q}$ in a color octet state. A gluon with virtuality $Q^2 = \mathcal{O}(M_B^2)$ is needed to equilibrate the large momentum fraction carried by the b quark in the $\overline{B}$ wavefunction. The amplitude then factors into a hard QCD/electroweak subprocess amplitude for quarks which are collinear with their respective hadrons: $T_H([b(x)\overline{u}(1-x)] \to [q(y)\overline{u}(1-y)]_1[q(z)\overline{q}(1-z)]_2)$ convoluted with the distribution amplitudes $\phi(x,Q)$ [3] of the incident and final hadrons:

$$\mathcal{M}_{octet}(B \to M_1 M_2) = \int_0^1 dz \int_0^1 dy \int_0^1 dx$$

$$\phi_B(x,Q) T_H(x,y,z) \phi_{M_1}(y,Q) \phi_{M_2}(z,Q).$$

Here $x = k^+/p_H^+ = (k^0+k^z)/(p_H^0+p_H^z)$ are the light-cone momentum fractions carried by the valence quarks.

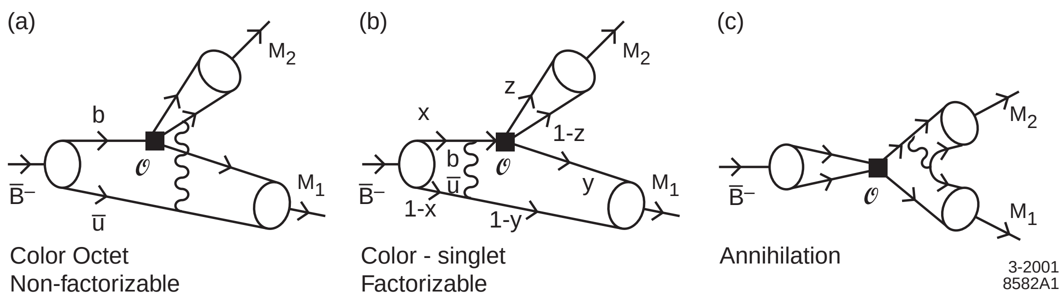

Figure 1: Three representative contributions to exclusive B decays to meson pairs in PQCD. The operators $\mathcal{O}$ represent the QCD-improved effective weak interaction

There are a several features of QCD which are required to ensure the consistency of the PQCD approach: (a) the effective QCD coupling $\alpha_s(Q^2)$ needs to be under control at the relevant scales of B decay; (b) the distribution amplitudes of the hadrons need to satisfy convergence properties at the endpoints; and (c) one requires the coherent cancelation of the couplings of soft gluons to color-singlet states. This property, color transparency [4], is a fundamental coherence property of gauge theory and leads to diminished final-state interactions and corrections to the PQCD factorizable contributions. Color transparency is not in contradiction with Watson's theorem for the factorization of final state phases since the energy in B decay is well above the regime where the scattering of the final states is elastic [5].

An innovative experiment (E791) at Fermilab has recently reported a direct determination of the quark-antiquark pion light-cone wavefunction by measuring the momentum distribution of diffractive dijet production on nuclei, $\pi A \to$ jet jet A, where the final nucleus stays intact [6]. The measured pion wavefunction appears to be similar to the form $\phi_M(x) \propto x(1-x)$ which is the asymptotic solution to the evolution equation for the pion distribution amplitude.

The E791 experiment [7] has also provided a remarkable confirmation of color transparency. Color transparency implies that when a fast hadron fluctuates into configurations of small transverse size, it can interact weakly and thus transverse a nucleus with minimal interactions [8, 9]. E791 finds that the diffractive dijet production process occurs coherently on each nucleon in the nucleus without nuclear absorption when the jet transverse momentum exceeds 2 GeV/c. The EVA spectrometer experiment E850 [10] at Brookhaven has also reported striking effects of color transparency in quasi-elastic proton-proton scattering on nuclei.

The distribution amplitudes $\phi_H(x, Q)$, which control hard exclusive B decays are fundamental gauge-invariant hadronic wavefunctions [3]. These amplitudes describe the light-cone momentum distributions of the valence quarks with relative transverse momentum $k_\perp < Q$. Evolution equations can be derived within PQCD which describe the change of $\phi(x, Q)$ with respect to $\ell n Q^2$. They can also be expanded in terms of specific polynomial eigenfunctions determined by conformal symmetry and corresponding to local operators with given anomalous dimensions [11, 12]. The distribution amplitude can also be understood as the transverse momentum integral of the lowest Fock state of the light-cone wavefunction of the hadron $\psi(x_i, k_{\perp i})$, or equivalently it can be computed from the Bethe-Salpeter amplitude evaluated at equal light-cone time $\tau = t + z/c$ in light-cone gauge $A^+ = 0$. Typically, one is interested in the leading twist-2 amplitude where the helicities of the pseudoscalar or vector

mesons have opposite light-cone helicity. Since they are process independent, the distribution amplitudes can be measured in other high momentum exclusive reactions, such as $\gamma\gamma \to M\overline{M}$. See Fig. 2. One can also make a connection between exclusive $B \to M_1 M_2$ amplitudes with the di-meson distribution amplitudes of meson pairs determined by measurements of $\gamma^*\gamma \to \pi^+\pi^-$ [13].

The endpoint-regions of integration $x \to 1$, $y \to 1$, $z \to 1$ of the octet amplitude converge because of the sufficiently fast fall off of the distribution amplitudes. This follows from the form of the solutions of the QCD evolution equation and also from the requirement that the light-cone kinetic energy of a quark $(k_\perp^2 + m_\perp^2)/k^+$ has finite expectation value. Note also that $x \to 1$ or $x \to 0$ corresponds kinematically to $k_z \to -\infty$ in the rest frame wavefunction. The contributions of higher Fock states containing extra gluons or sea quarks are suppressed by powers of m_b^{-1} in light-cone gauge. Soft gluon corrections such as those leading to final-state interactions are suppressed by the color neutrality of the final state hadrons as they emerge with small color dipole moments. Insertions of collinear gluons from higher Fock states into the hard amplitude are power-law suppressed in light-cone gauge [14]. The proofs of factorization for exclusive B decays are generalizations of the analyses of exclusive amplitudes involving large momentum transfer [3, 15].

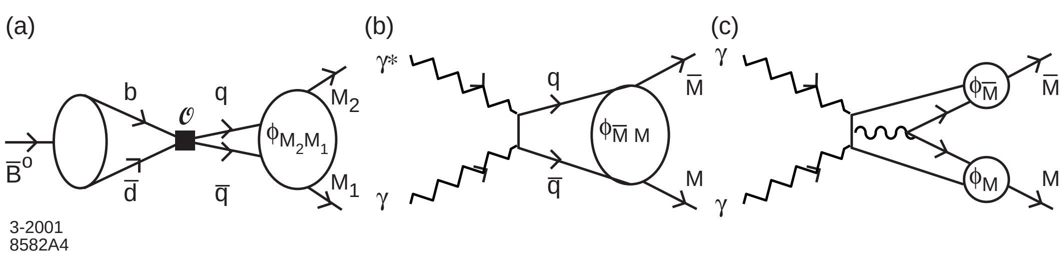

Figure 2: Similarities between exclusive $B \to M_1 M_2$ decays and $\gamma\gamma \to M\overline{M}$ decays in PQCD.

Can one trust the applicability of leading twist PQCD factorization theorems at momentum transfers involved in B decays? The CLEO collaboration [16] has verified the scaling and angular predictions for hard exclusive two-photon processes such as $\gamma^*\gamma \to \pi^0$ and $\gamma\gamma \to \pi^+\pi^-$ at momentum transfers comparable to the momenta that occur in exclusive B-decays. The CLEO data for the sum of $\gamma\gamma \to \pi^+\pi^+ + K^+K^-$ channels at $W = \sqrt{s} > 2.5$ GeV are in striking agreement with the perturbative QCD predictions. Moreover, the observed angular distribution shows a striking transition to the predicted QCD form as W is raised. The L3 experiment at LEP at CERN

[17] has also measured a number of exclusive hadron production channels in two-photon processes, providing important constraints on baryon and meson distribution amplitudes and checks of perturbative QCD factorization. The $\gamma^*\gamma \to \pi^0$ results are in close agreement with the scaling and normalization of the PQCD prediction, provided that the pion distribution amplitude $\phi_\pi(x, Q)$ is close to the $x(1-x)$ form, the asymptotic solution to the evolution equation, in agreement with dijet diffraction measurements from E791. It will be important to have measurements of processes such as $\gamma\gamma \to \pi^0\pi^0$ and $\gamma\gamma \to \rho^+\rho^-$ since they are particularly sensitive to the shape of meson distribution amplitudes [18].

The scale of the coupling which controls T_H in hard exclusive processes can be obtained using the BLM method; the scale is of order of the gluon virtuality and the perturbative coefficients are identical to those of conformal theory [19, 20]. The empirical success of dimensional counting rules for exclusive two-photon reactions and other exclusive amplitudes at moderate momentum transfer suggests that the QCD coupling $\alpha_s(Q^2)$ is effectively constant at low Q^2 [21].The effective charge $\alpha_\tau(s)$ measured in hadronic τ decays integrated up to invariant mass squared s has been shown to have this property [22]. All of these aspects of QCD are essential for the analysis of exclusive B decays.

In contrast to the color-octet contribution, the contribution of Fig. 1(b) to the $B \to M_1 M_2$ amplitude where the light quarks of meson M_2 are produced in a color-singlet state has an apparent $(1-y)^{-2}$ endpoint singularity in the limit of large heavy quark mass [23]. Thus one cannot guarantee the convergence of the convolution of T_H solely from the fall-off of the distribution amplitudes; thus strictly speaking such amplitudes cannot be analyzed at fixed order in PQCD. The simplest approach, taken by Buchalla *et al.*, is to identify the required color-singlet transition amplitude $B \to M_1$ with the transition form factor $F_{B\to M_1 \ell\overline{\nu}}(M_2^2)$ measurable in semileptonic decays.

Despite the complication from the endpoint sensitivity of the color-singlet contributions, some predictability is still possible from PQCD. Li *et al.* [24, 25] have noted that Sudakov form factors control will effectively suppress any power-law endpoint singularities which appear when the quarks have finite transverse momentum. This type of Sudakov suppression is also an essential element in the analysis of the $x \to 1$ endpoint, finite k_T, contributions to the elastic form factors of nucleons [3]. Physically, such Sudakov effects occur when a quark near its mass shell is forced to decelerate without gluonic radiation. The resummation of a double-logarithmic series

gives the form ($C_F = \frac{N_C^2-1}{2N_C}$): $S(Q^2) = \exp -\frac{\alpha_s C_F}{4\pi}\ell n^2 \frac{Q^2}{k_\perp^2}$ at fixed QCD coupling which suppresses the long-distance contributions in the large b region. The exponent becomes a $\ell n \times \ell n \ell n$ form when the running of the QCD coupling is taken into account. Keum and Li [25] find that almost all of the contribution to the exclusive matrix elements then come from the integration region where $\alpha_s/\pi < 0.3$ and the perturbative QCD treatment can be is justified. Thus for PQCD contributions which are Sudakov form factor controlled, the effective scale of the subprocess is $Q^2 = \mathcal{O}(m_b \Lambda_{QCD})$. Li *et al.* argue that the presence of Sudakov suppression allows the applicability of PQCD to exclusive decays.

The Sudakov suppression approach also allows new predictions for the phase structure of the annihilation graphs such as Fig. 1(c). In this case, the endpoint behavior of the higher twist wavefunctions which appear in the helicity parallel amplitudes as determined by QCD evolution may not be sufficient to control the endpoint singularity region. Such amplitudes, however, would be controlled by the Sudakov form factors. Since the scale of the exchanged gluon momentum is of order $\Lambda_{QCD} m_b$, the coupling strength of the exchanged gluons would be enhanced by a larger effective $\alpha_s(Q^2)$, thus leading to the possibility of enhanced CP-violating phases [25]. Further discussion of the phenomenology of exclusive QCD decays based on the factorization theorems and the Sudakov-extended analyses has been given by Li [24], Neubert [26], and Becher *et al.* [27] It should also be possible to extend the PQCD formalism to the exclusive decays of B mesons to baryon pairs.

2 Exact Representation of Electroweak Matrix Elements in terms of Light-Cone Wavefunctions

The natural calculus for describing the bound-state structure of relativistic composite systems in quantum field theory is the light-front Fock expansion which encodes the properties of a hadrons in terms of a set of frame-independent n-particle wavefunctions. A remarkable advantage of the light-cone formalism is that exclusive semileptonic B-decay amplitudes such as $B \to A\ell\overline{\nu}$ can be evaluated *exactly* [28]. The time-like decay matrix elements require the computation of the diagonal matrix element $n \to n$ [See Fig. 3(a)] where parton number is conserved, and the off-diagonal $n+1 \to n-1$ convolution where the current operator annihilates a $q\overline{q'}$ pair in the initial B wavefunction, as illustrated in Fig. 3(b). This term is a consequence of the fact that the time-like decay $q^2 = (p_\ell + p_{\overline{\nu}})^2 > 0$ requires a positive light-cone momentum

fraction $q^+ > 0$. The sum of amplitudes is required to obtain a covariant result. A similar form also controls deeply virtual Compton scattering. In contrast, for space-like currents, one can choose $q^+ = 0$, as in the Drell-Yan-West representation of the space-like electromagnetic form factors.

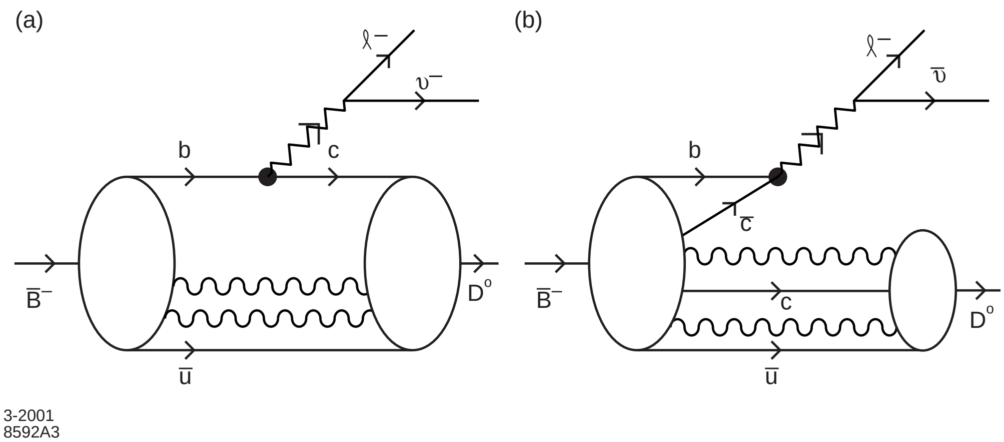

Figure 3: Diagonal and off-diagonal contributions to time-like B decays. Both contributions are required by Lorentz covariance.

The light-cone Fock expansion is defined as the projection of an exact eigensolution of the full light-cone quantized Hamiltonian on the solutions of the free Hamiltonian with the same global quantum numbers. The coefficients of the Fock expansion are the complete set of n-particle light-cone wavefunctions, $\{\psi_n(x_i, k_{\perp i}, \lambda_i)\}$. The coordinates $x_i, k_{\perp i}$ are internal relative coordinates, independent of the total momentum of the bound state, and satisfy $0 < x_i < 1$, $\sum_i^n x_i = 1$ and $\sum_i^n k_{\perp i} = 0_\perp$. The light-cone wavefunctions are thus independent of the total hadronic momentum.

The evaluation of the semileptonic decay amplitude $B \to A\ell\overline{\nu}$ requires the matrix element of the weak current between hadron states $\langle A|j^\mu(0)|B\rangle$. (See Fig. 4.) We can choose the Lorentz frame where the outgoing leptonic current carries $q^\mu = (q^+, q_\perp, q^-) = \left(\Delta P^+, q_\perp, \frac{q^2+q_\perp^2}{\Delta P^+}\right)$. In the limit $\Delta \to 0$, the matrix element for the $+$ vector current coincides with the Drell-Yan formula for spacelike form factors.

The diagonal (parton-number-conserving) matrix element has the form of a convolution of $\psi^\dagger_{A(n)}(x'_i, \vec{k}'_{\perp i}, \lambda_i)$ and $\psi_{B(n)}(x_i, \vec{k}_{\perp i}, \lambda_i)$ where $x'_1 = \frac{x_1 - \Delta}{1-\Delta}$, $\vec{k}'_{\perp 1} = \vec{k}_{\perp 1} - \frac{1-x_1}{1-\Delta}\vec{q}_\perp$ for the struck quark and $x'_i = \frac{x_i}{1-\Delta}$, $\vec{k}'_{\perp i} = \vec{k}_{\perp i} + \frac{x_i}{1-\Delta}\vec{q}_\perp$ for the $(n-1)$ spectators. If quark masses are neglected the vector and axial currents conserve helicity. Note that $\sum_i^n x'_i = 1$, and $\sum_i^n \vec{k}'_{\perp i} = \vec{0}_\perp$.

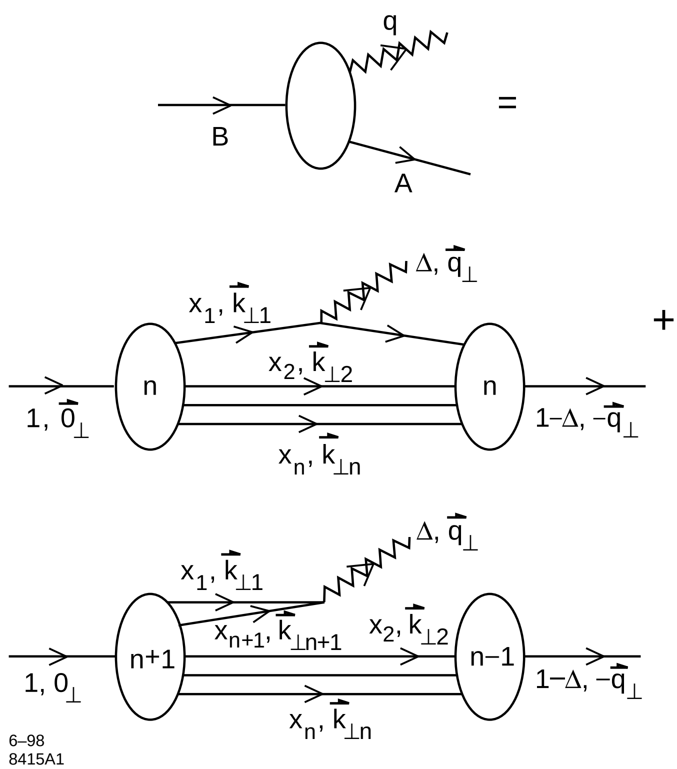

Figure 4: Exact representation of electroweak decays and time-like form factors in the light-cone Fock representation.

For the $n+1 \to n-1$ off-diagonal term, consider the case where partons 1 and $n+1$ of the initial wavefunction annihilate into the leptonic current leaving $n-1$ spectators. Then $x_{n+1} = \Delta - x_1$, $\vec{k}_{\perp n+1} = \vec{q}_\perp - \vec{k}_{\perp 1}$. The remaining $n-1$ partons have total momentum $((1-\Delta)P^+, -\vec{q}_\perp)$. The final wavefunction then has arguments $x'_i = \frac{x_i}{(1-\Delta)}$ and $\vec{k}'_{\perp i} = \vec{k}_{\perp i} + \frac{x_i}{1-\Delta}\vec{q}_\perp$. The arguments of the final-state wavefunction satisfy $\sum_{i=2}^{n} x'_i = 1$, $\sum_{i=2}^{n} \vec{k}'_{\perp i} = \vec{0}_\perp$.

The off-diagonal $n+1 \to n-1$ contributions give a new perspective for the physics of B-decays. A semileptonic decay involves not only matrix elements where a quark changes flavor, but also a contribution where the leptonic pair is created from the annihilation of a $q\overline{q'}$ pair within the Fock states of the initial B wavefunction. The semileptonic decay thus can occur from the annihilation of a nonvalence quark-antiquark pair in the initial hadron. This feature will carry over to exclusive hadronic B-decays, such as $B^0 \to \pi^- D^+$. In this case the pion can be produced from the coalescence of a $d\overline{u}$ pair emerging from the initial higher particle number Fock wavefunction of the B. The D meson is then formed from the remaining quarks after

the internal exchange of a W boson.

In principal, the evaluation of the hadronic matrix elements needed for B-decays and other exclusive electroweak decay amplitudes requires knowledge of all of the light-cone Fock wavefunctions of the initial and final state hadrons. In practice only the first few Fock components have significant magnitude. In the case of model gauge theories such as QCD(1+1) [29] or collinear QCD [30] in one-space and one-time dimensions, the complete evaluation of the light-cone wavefunction is possible for each baryon or meson bound-state using the DLCQ method. In fact, solutions of the t'Hooft model have been used to model physical B-decays [31, 32, 33, 34].

3 Canonical Form for Light-Cone Wavefunctions at Small Transverse Momenta

One can understand some basic features of the LC wavefunctions of hadrons, including the wavefunction of the B meson itself by considering its canonical form at low transverse moments $k_{\perp i}$: [35]

$$\psi_n(x_i, k_{\perp i}) = \frac{\Gamma(x_i, k_{\perp i})}{[\delta^2 + \sum_{j=1}^{n} \frac{(x_j - \hat{x}_j)^2}{x_j}]}$$

where $\delta^2 = 2B.E./M_H$, $\hat{x}_j = m_{\perp j}/\sum_{k=1}^{n} m_{\perp k}$, and the transverse mass squared is $m_{\perp j}^2 = m_j^2 + k_{\perp j}^2$. Here Γ is the convolution of the interaction kernel V_{LC} with the light-cone wavefunction and can be considered as slowly varying compared to the peaking of the denominator. The spin structure of the numerator Γ can be largely constructed using light-cone angular momentum conservation [36]. The canonical form follows directly from the bound-state light-cone Hamiltonian equation of motion $[M^2 - \sum_{i=1}^{n} m_\perp^2]\,\psi = V_{LC}\Psi$. One sees that at small values of $k_{\perp i}$, the light-cone wavefunctions peak at value $x_i = \hat{x}_i$, the equal rapidity, minimal off-shell configuration. The spread of the wavefunction where the wavefunction falls to half of its value is $< (x_i - \hat{x}_i)^2 >= x_i \delta^2$ and is largest for the heaviest partons. The full wavefunction can then be constructed from this low $k_\perp$ starting point by iterating the equation of motion [3].

DLCQ (discretized light-cone quantization) is a method which solves quantum field theory by direct diagonalization of the light-front Hamiltonian [37, 38]. The DLCQ method has had much success in solving quantum field theories in low space-time dimensions and has also found utility in string theory. There has been progress

recently in systematically developing the computation and renormalization methods needed to make DLCQ viable for QCD in physical spacetime. There has also been considerable success in applying these methods to simple quantum field theories in 3+1 dimensions [39]. The distribution amplitude of the pion has recently been computed using a combination of the discretized light-cone quantization and transverse lattice methods. [40]

4 Intrinsic Charm and B Decay

The complete Fock state description of the B meson wavefunction contains higher particle number states such as $|b\overline{u}g>$, $|b\overline{u}s\overline{s}>$, and $|b\overline{u}c\overline{c}>$. Such Fock components arise from gluon splitting and from diagrams in which the sea quarks are multi-connected to the valence quarks. Since the light-cone wavefunction peaks at $x_i \propto m_{i\perp}$, the intrinsic charm quarks are found at relatively high x [41]. Evidence for a 1% probability of intrinsic charm in the proton has been given by Harris *et al.* [42] based on an analysis of the EMC measurement of the proton's charm structure function at large x. Intrinsic charm can also explain the $J/\psi \to \rho\pi$ puzzle [43]. Franz *et al.* [44] have recently given a rigorous operator product expansion of intrinsic heavy quarks showing that the momentum fraction carried by intrinsic heavy quarks , $x_{Q\overline{Q}}$ scales as p_H^2/m_Q^2 where p_H is the characteristic internal momentum in the hadron. This is strictly a non-Abelian effect and only occurs for intrinsic charm pairs in a color octet configuration [45]. Since the B wavefunction is more compact compared to light hadrons (a reduced mass effect), the magnitude intrinsic charm is dynamically enhanced in the B wavefunction.

The presence of intrinsic charm quarks in the B wavefunction provides new mechanisms for B decays. For example, Chang and Hou [46] have considered the production of final states with three charmed quarks such as $B \to J/\psi D\pi$ and $B \to J/\psi D^*$ which arises naturally when the b quark of the intrinsic charm Fock state $|b\overline{u}c\overline{c}>$ decays $b \to c\overline{u}d$. See Fig. 5(a). In fact, the J/ψ spectrum for inclusive $B \to J/\psi X$ decays measured by CLEO and Belle does show a distinct enhancement at low J/ψ momentum where such decays would occur [46]. Another interesting possibility is that the enhancement at low J/ψ momentum is due to final states containing baryon pairs. $B \to J/\psi\overline{p}\Lambda$ distribution which would also allow the study of di-baryons or the bound states of baryons with the J/ψ which can arise via the QCD van der Waals potential at low relative velocity [47].

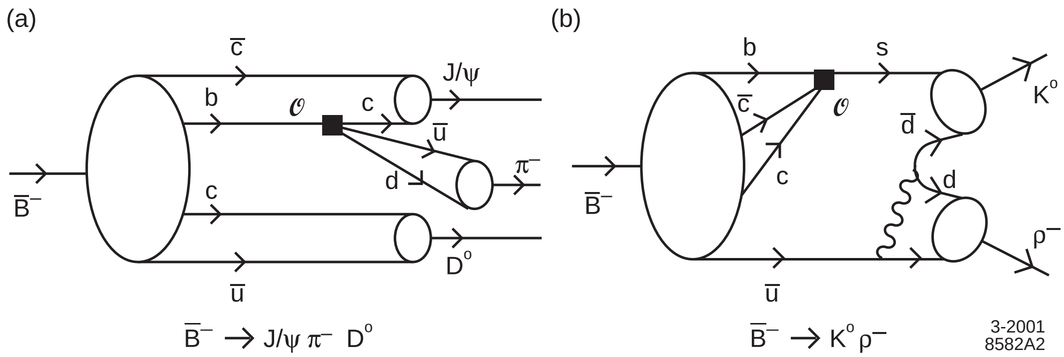

Figure 5: Examples of intrinsic charm effects in exclusive B decays. In (a) the intrinsic charm of the B wavefunction leads to final states with three charmed quarks. In (b), the annihilation of the intrinsic charm quarks in the effective weak interaction allows a decay to proceed without CKM suppression.

The existence of intrinsic charm in the B wavefunction, even at a few percent probability, allows exclusive decays of the B meson which evade the hierarchy of the CKM matrix [45]. This can lead to significant modifications of standard predictions for exclusive decays such as $B \to \rho K$. In the standard analyses, the tree-level contributions involving $b \to u\overline{u}s$ are CKM suppressed, whereas the penguin contributions which are not CKM suppressed are numerically suppressed by Wilson coefficients.

An example of an intrinsic charm contribution to $B \to \rho\pi$ is illustrated in Fig. 5(b). The three-particle weak annihilation process $bc\overline{c} \to sg$ can occur via W^- exchange tree-level operators using the large V_{cb} and V_{cs} matrix element. The non-valence quark annihilation process thus provides a new mechanism for decays such as $B^- \to \rho^0 K^-$. The relatively small intrinsic charm probability of $\sim 4\%$ can thus be offset by the comparatively large CKM matrix elements of the effective weak interaction. Furthermore, since the same initial and final states are involved, the intrinsic charm annihilation amplitudes can interfere with the conventional amplitudes, leading to a significant correction to the traditional analyses for decay amplitudes involving the intrinsic charm Fock state.

5 Conclusions

Light-cone wavefunctions provide a fundamental frame-independent description of hadrons in terms of their quark and gluon degrees of freedom at the amplitude level and a natural description of the exclusive amplitudes which occur in B decay. For

example, the semi-leptonic decay amplitudes of hadrons can be computed exactly in terms of the diagonal $\Delta n = 0$ and off-diagonal $\Delta n = 2$ overlap of the hadronic light-cone wavefunctions. Since the momentum transfers are large, exclusive B decays into hadron pairs can be factorized as convolutions of hard scattering amplitudes involving the weak interaction together with universal hadron distribution amplitudes, thus providing a new perturbative QCD-based phenomenology. These factorization theorems, together with the universality of the distribution amplitudes, provide a profound connection between hard scattering exclusive processes such as elastic form factors, two-photon reactions, and heavy hadron decays. We can understand the shape of the light-cone wavefunctions using a canonical form which is valid at low values of the transverse momenta. Color transparency implies minimal phases from final-state interactions, a fact which is critical to the interpretation of CP violation in B physics. I have also emphasized the fact the existence of intrinsic charm Fock states in the B-meson wavefunction can enhance the production of final states of the B with three charmed quarks, such as $B \to J/\psi D\pi$, as well as lead to the breakdown of the CKM hierarchy.

6 Acknowledgments

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References

[1] M. Beneke, G. Buchalla, M. Neubert and C. T. Sachrajda, *Nucl. Phys.* **B591**, 313 (2000); [hep-ph/0006124].

[2] Y. Y. Keum, H. Li and A. I. Sanda, hep-ph/0004173.

[3] G. P. Lepage and S. J. Brodsky, *Phys. Rev.* **D22**, 2157 (1980).

[4] S. J. Brodsky and A. H. Mueller, *Phys. Lett.* **B206**, 685 (1988).

[5] M. Suzuki and L. Wolfenstein, *Phys. Rev.* **D60**, 074019 (1999) [hep-ph/9903477].

[6] E. M. Aitala *et al.* [Fermilab E791 Collaboration], hep-ex/0010043.

[7] E. M. Aitala *et al.* [Fermilab E791 Collaboration], hep-ex/0010044.

[8] G. Bertsch, S. J. Brodsky, A. S. Goldhaber and J. F. Gunion, *Phys. Rev. Lett.* **47**, 297 (1981).

[9] L. Frankfurt, G. A. Miller and M. Strikman, hep-ph/0010297.

[10] A. Leksanov *et al.* [E850 Collaboration], nucl-ex/0009008.

[11] S. J. Brodsky, Y. Frishman, G. P. Lepage and C. Sachrajda, *Phys. Lett.* **B91**, 239 (1980).

[12] V. M. Braun and I. E. Filyanov, *Z. Phys.* **C48**, 239 (1990).

[13] M. Diehl, T. Gousset and B. Pire, *Phys. Rev.* **D62**, 073014 (2000) [hep-ph/0003233].

[14] P. P. Srivastava and S. J. Brodsky, hep-ph/0011372.

[15] V. L. Chernyak and A. R. Zhitnitsky, *Phys. Rept.* **112**, 173 (1984).

[16] J. Dominick *et al.* [CLEO Collaboration], *Phys. Rev.* **D50**, 3027 (1994) [hep-ph/9403379]. J. Gronberg *et al.* [CLEO Collaboration], *Phys. Rev.* **D57**, 33 (1998) [hep-ex/9707031]; and H. Paar, presented at Photon 2000.

[17] H. Vogt [L3 Collaboration], *Nucl. Phys. Proc. Suppl.* **79**, 283 (1999).

[18] For a recent review, see S. J. Brodsky, hep-ph/0010176. S. J. Brodsky and G. P. Lepage, *Phys. Rev.* **D24**, 1808 (1981).

[19] S. J. Brodsky, G. P. Lepage and P. B. Mackenzie, *Phys. Rev.* **D28**, 228 (1983).

[20] S. J. Brodsky, E. Gardi, G. Grunberg and J. Rathsman, hep-ph/0002065.

[21] S. J. Brodsky, C. Ji, A. Pang and D. G. Robertson, *Phys. Rev.* **D57**, 245 (1998) [hep-ph/9705221].

[22] S. Brodsky, S. Menke, and C. Merino, in preparation.

[23] A. Szczepaniak, E. M. Henley and S. J. Brodsky, *Phys. Lett.* **B243**, 287 (1990).

[24] H. Li, hep-ph/0102013.

[25] Y. Keum and H. Li, *Phys. Rev.* **D63**, 074006 (2001) [hep-ph/0006001].

[26] M. Neubert, hep-ph/0012204.

[27] T. Becher, M. Neubert and B. D. Pecjak, hep-ph/0102219.

[28] S. J. Brodsky and D. S. Hwang, *Nucl. Phys.* **B543**, 239 (1999); [hep-ph/9806358]. C. Ji and H. Choi, hep-ph/0009281.

[29] K. Hornbostel, S. J. Brodsky, and H. C. Pauli, *Phys. Rev.* **D41** 3814 (1990).

[30] F. Antonuccio and S. Dally, *Phys. Lett.* **B348**, 55 (1995); *Phys. Lett.* **B376**, 154 (1996); *Nucl. Phys.* **B461**, 275 (1996).

[31] B. Grinstein and R. F. Lebed, *Phys. Rev.* **D59**, 054022 (1999) [hep-ph/9805404].

[32] R. F. Lebed and N. G. Uraltsev, *Phys. Rev.* **D62**, 094011 (2000) [hep-ph/0006346].

[33] I. Bigi, M. Shifman, N. Uraltsev and A. Vainshtein, *Phys. Rev.* **D59**, 054011 (1999) [hep-ph/9805241].

[34] M. Neubert, hep-ph/0006265.

[35] S. J. Brodsky and D. S. Hwang, in preparation.

[36] S. J. Brodsky, D. S. Hwang, B. Ma and I. Schmidt, *Nucl. Phys.* **B593**, 311 (2001) [hep-th/0003082].

[37] H. C. Pauli and S. J. Brodsky, *Phys. Rev.* **D32**, 2001 (1985).

[38] S. J. Brodsky, H. Pauli and S. S. Pinsky, *Phys. Rept.* **301**, 299 (1998) [hep-ph/9705477].

[39] S. J. Brodsky, J. R. Hiller and G. McCartor, *Phys. Rev.* **D60**, 054506 (1999) [hep-ph/9903388].

[40] S. Dalley, hep-ph/0007081. S. K. Seal and M. Burkardt, *Nucl. Phys. Proc. Suppl.* **90**, 233 (2000).

[41] S. J. Brodsky, P. Hoyer, C. Peterson and N. Sakai, *Phys. Lett.* **B93**, 451 (1980).

[42] B. W. Harris, J. Smith and R. Vogt, *Nucl. Phys.* **B461**, 181 (1996) [hep-ph/9508403].

[43] S. J. Brodsky and M. Karliner, *Phys. Rev. Lett.* **78**, 4682 (1997) [hep-ph/9704379].

[44] M. Franz, M. V. Polyakov and K. Goeke, *Phys. Rev.* **D62**, 074024 (2000) [hep-ph/0002240].

[45] S. J. Brodsky and S. Gardner, in preparation.

[46] C. V. Chang and W. Hou, hep-ph/0101162.

[47] S. J. Brodsky and F. S. Navarra, *Phys. Lett.* **B411**, 152 (1997) [hep-ph/9704348].